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Two Higgs Doublets from Fermion Condensation

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Abstract: We consider the most generic situation in models where the electroweak symmetry is broken by the condensation of a strongly coupled fermion sector, such as for instance a fourth generation. We study the scalar content resulting from the condensation of both the up and the down type fermions, corresponding to a two-Higgs doublet model. We estimate the scalar spectrum using the Nambu–Jona-Lasinio model, improved by the renormalization group. We show that the scalar spectrum is generically lighter than for the case with only one right-handed fermion condensing and that due to a remnant Peccei-Quinn symmetry the lightest state is the pseudo-scalar, with masses ranging typically from 10 GeV to 120 GeV. We discuss the phenomenological consequences of this distinct spectrum.

Keywords: electroweak symmetry breaking; dynamical higgs mechanism; two Higgs doublet models

Contents

1 Introduction

The standard model (SM) gives a very successful description of the gauge interactions of all elementary particles known up to now [1]. On the other hand, the origin of the electroweak symmetry breaking (EWSB) sector, responsible for all mass scales, remains an open question. An alternative to the elementary scalar doublet introduced in the SM, is the possibility that a new interaction dynamically generates a new scale at low energies. This paradigm, which is at work in the way the strong interactions generate the hadronic scale, is used in technicolor theories[2, 3], where the new strong interaction resulting in EWSB is confining. This results in a spectrum of "technihadrons" formed by the strongly coupled fermions. Alternatively, we could consider the possibility that the new strong interaction is spontaneously broken at around the scale where it becomes supercritical. In this case, the condensing fermions would not be confined. This raises the possibility of using SM fermions to break the electroweak symmetry through their condensation [4]. However, the dynamical mass acquired by the condensing fermions must be approximately $(500 - 700)$ GeV if we want the scale at which condensation takes place to be close to $O(1)$ TeV. Thus, we must introduce new fermions [5] in addition to the new strong interaction. This causes the condensates and leads to large dynamical masses for the condensing fermions, as well as to EWSB.

In order to fully realize this scenario, various ingredients must be in place. In addition to the new fermions, a new strong interaction spontaneously broken at the TeV scale and strongly coupled to them must be present. Furthermore, higher-dimensional operators involving the massless fermions with the condensing ones must be present at low energies in order to feed masses and mixing to the SM fermions. A complete model was proposed in Ref.[6] based on the SM with four generations in the bulk of a compact extra dimension

with an AdS background. We will refer to this as the flavor AdS_5 model. This model, which minimally extends bulk Randall-Sundrum models [7, 8] by adding a highly localized fourth generation in the bulk, constitutes an existence proof as well as a source of detailed predictions [9, 10].

In Ref.[6] it was assumed for simplicity that only one of the fourth-generation quarks condenses. This leads to a composite scalar sector that corresponds to that of one scalar doublet, just as in the SM. The Higgs mass in that case was at the edge of perturbative unitarity. However, it is more natural to assume that both up and down-type fermions form condensates since the new strong interaction should largely respect isospin symmetry in order to avoid T parameter constraints. If both up and down-type fermions condense, the resulting composite scalar spectrum corresponds to that of a two-Higgs doublet model [11]. In this paper we will study this possibility, particularly the low energy scalar theory resulting from double condensation. Although we are taking inspiration from the flavor AdS_5 model to generate the necessary interaction among fourth-generation quarks, our results are valid any ultra-violet (UV) completion of the four-fermion interactions giving rise to the condensation. Thus, the resulting low energy scalar spectrum and phenomenology should correspond to a large class of theories, independently of whether the condensing fermions are fourth generation quarks or just generic new fermions strongly coupled to a spontaneously broken gauge interaction.

The rest of the paper is organized as follows: in Section 2 we specify the new fermion interactions introduced the couplings to massive gauge bosons. Several properties of the resulting scalar spectrum at energies below the gauge boson mass are specified already here. In Section 3 we compute the low energy scalar spectrum. We do this by solving the renormalization group equations (RGE) for the scalar self-couplings and the fourth-generation Yukawa couplings. In Section 4 we study the electroweak precision constraints in a model with four fermion generations and the scalar spectrum obtained in Section 3. The constraints from flavor physics and direct searches are discussed in Section 5, whereas we conclude in Section 6.

2 A Two-Higgs Doublet from Fermion Condensation

We consider a scenario with a set of new fermions coupled to a new strong interaction. This is mediated by massive gauge bosons, possibly remnant from the spontaneous breaking of a gauge symmetry at the TeV scale. For instance, in the flavor AdS_5 models mentioned in the previous section, the strongest interaction among the fourth-generation quarks is provided by the KK gluons and the fermions are the zero modes of fourth-generation quarks. More generally, we can consider that there is a new interaction spontaneously broken above the TeV scale, which will be responsible for fermion condensation. Thus, the minimal fermion content we consider is given by

$$Q^a = \begin{pmatrix} U^a \\ D^a \end{pmatrix}_L , \quad U^a, \, D^a \ , \tag{2.1}$$

where the index a refers to the new interaction felt by the $SU(2)_L$ doublet Q as well as by the singlets U and D. The fermions in (2.1) are the ones that will condense. Additional matter content might be necessary in order to cancel anomalies. For instance, if the fermions in (2.1) are fourth-generation quarks, leptons must be added in order to get an anomaly free theory. Just as in this case, we assume that in general the additional fermions needed to cancel anomalies will not participate in condensation. The new interaction could also couple to the lighter quarks, although more weakly. Here we will not consider these interactions for simplicity. We further assume that the new strong interaction is spontaneously broken and that massive gauge bosons are integrated out to lead to four-fermion interactions of the form

$$\begin{aligned}\mathcal{L}_{\text{eff}} &= \frac{g_L g_u}{M_G^2}\,\bar{Q}\gamma_\mu T^a Q\,\bar{U}\gamma^\mu T^a U \\ &+ \frac{g_L g_d}{M_G^2}\,\bar{Q}\gamma_\mu T^a Q\,\bar{D}\gamma^\mu T^a D\ .\end{aligned} \tag{2.2}$$

where T^a are the symmetry generators, and g_L, g_u and g_d are the gauge couplings of Q, U and D, respectively. For instance, in the case where the new fermions are fourth-generation quarks, the T^a are the $SU(3)_c$ generators, and the massive gauge boson is a color-octect, just as in the AdS$_5$ flavor model of Ref. [6]. Here we will consider a generic case, since the nature of the new interaction as well as the identity of the condensing fermions is not crucial in determining the low-lying scalar spectrum, as we will see below.

The four-fermion interactions of interest leading to EWSB are

$$\mathcal{L}_{\text{eff.}} = G_U\bar{Q}U\bar{U}Q + G_D\bar{Q}D\bar{D}Q \tag{2.3}$$

where the Nambu–Jona-Lasinio (NJL) couplings are just

$$G_U \equiv \frac{g_L g_u}{M_G^2}, \qquad G_D \equiv \frac{g_L g_d}{M_G^2}\ . \tag{2.4}$$

In order to generate non-trivial values for the quark condensates $\langle\bar{Q}U\rangle$ and $\langle\bar{Q}D\rangle$, the couplings G_U and/or G_D must be above some critical value. In Ref. [6] only one of the couplings (G_U) was considered supercritical, leading to a composite scalar spectrum consisting of one Higgs doublet. Here we consider a more general case, with both couplings supercritical. As mentioned earlier, this results in a two-Higgs doublet composite scalar spectrum, with very distinct masses. This scenario will turn out to be quite different from the one considered in [6], even if we just consider the neutral scalar masses.

The four-fermion interactions can be bosonized by introducing the auxiliary fields $\tilde{\Phi}_U$ and Φ_D, which are $SU(2)_L$ doublets with hypercharges $h_U = -1/2$, and $h_D = 1/2$. Then the NJL interactions of (2.3) can be shown to be equivalent to

$$\mathcal{L}_{\text{eff.}} = Y_U(\bar{Q}\tilde{\Phi}_U U + \text{h.c.}) + Y_D(\bar{Q}\Phi_D D + \text{h.c.}) - M_G^2\Phi_U^\dagger\Phi_U - M_G^2\Phi_D^\dagger\Phi_D\ , \tag{2.5}$$

with the appropriate identifications:

$$Y_U^2 = g_L g_u, \qquad Y_D^2 = g_L g_d\ . \tag{2.6}$$

The auxiliary fields in (2.5) defined at the scale M_G develop kinetic terms and self-interactions below that scale

$$\mathcal{L}_{\text{kin.}} = Z_{\Phi_U}(D_\mu \Phi_U)^\dagger D^\mu \Phi_U + Z_{\Phi_D}(D_\mu \Phi_D)^\dagger D^\mu \Phi_D \ , \tag{2.7}$$

where the field renormalizations obey $Z_{\Phi_i}(\mu) \to 0$ as $\mu \to M_G$. This leads to the renormalization of the parameters of the potential. For instance, at leading order, the Yukawa couplings renormalize according to

$$Y_U \to \frac{Y_U}{\sqrt{Z_{\Phi_U}}}, \qquad Y_D \to \frac{Y_D}{\sqrt{Z_{\Phi_D}}},$$

The mass terms, on the other hand, will get corrections that are quadratic in the cutoff M_G, leading to new renormalized masses at the scale $\mu < M_G$:

$$\mu_U^2 = M_G^2 - \frac{g_L g_u N_g}{8\pi^2}\left(M_G^2 - \mu^2\right) \tag{2.8}$$

$$\mu_D^2 = M_G^2 - \frac{g_L g_d N_g}{8\pi^2}\left(M_G^2 - \mu^2\right) \ , \tag{2.9}$$

where N_g is the size of the fermion representation in the broken group. For instance, if the new fermions are fourth-generation quarks $N_g = N_c$, the number of colors. Finally, there will be scalar self-interactions generated by fermion loops, which also get renormalized by $Z_{\Phi_{U,D}}$. We will consider them in detail in the next section, where we study the scalar potential.

As we will see below, the minimization condition on the scalar potential will almost certainly require that, for $\mu \ll M_G$, both $m_U^2 < 0$ and $m_D^2 < 0$ be satisfied in order for Φ_U and Φ_D to get a VEV each, v_U and v_D respectively. In terms of the underlying gauge theory, this translates into a criticality condition for the gauge couplings

$$g_L g_u > \frac{8\pi^2}{N_g}, \qquad g_L g_d > \frac{8\pi^2}{N_g}, \tag{2.10}$$

which coincides with the conditions to make the NJL couplings in (2.3) supercritical to form both fermion condensates, $\langle \bar{Q} U\rangle$ and $\langle \bar{Q} D\rangle$.

We notice that the tree-level auxiliary lagrangian (2.5) *lacks* a term mixing Φ_U with Φ_D. This means that the theory is invariant under the Peccei-Quinn symmetry

$$\begin{array}{ccc} Q \to e^{-i\theta} Q & U \to e^{i\theta} U & D \to e^{i\theta} D \\ \Phi_U \to e^{2i\theta}\Phi_U & \Phi_D \to e^{-2i\theta}\Phi_D \ , & \end{array} \tag{2.11}$$

with θ an arbitrary phase. The mixing term in the bosonized theory has the form

$$\mathcal{L}_{\text{mix}} = \mu_{UD}^2(\Phi_U^\dagger \Phi_D + \text{h.c.}) \tag{2.12}$$

and, although is not generated by fermion loops perturbatively, it will be generated by instanton effects. As we will see in the next section, in the $\mu_{UD} = 0$ limit the pseudo-scalar state becomes massless. Instanton effects associated to the new strong interaction will

lift this "axion" mass, but it will remain the lightest state in the spectrum. Going back to fermion degrees of freedom, the term leading to the mixing written purely in terms of fermion fileds is

$$\mathcal{L}_{\rm mix} = G_{UD}(\bar{Q}D\bar{U}^c\tilde{Q} + {\rm h.c.}) \ , \tag{2.13}$$

with

$$\tilde{Q} \equiv -i\sigma_2 Q \ . \tag{2.14}$$

This term is generated by instantons and in fact is just the 't'Hooft flavor determinant [12]

$$\begin{aligned}\mathcal{L}_{inst.} &= \frac{\kappa}{M_G^2}\mathbf{det}\left[\bar{Q}_L Q_R\right] \\ &= \frac{\kappa}{M_G^2}\left[(\bar{D}_L D_R)(\bar{U}_L U_R) - (\bar{U}_L D_R)(\bar{D}_L U_R) + {\rm h.c.}\right]\end{aligned} \tag{2.15}$$

where κ is an $O(1)$ dimensionless coefficient depending on details of the instanton calculation. Thus, the mixing parameter in (2.13) $G_{UD} = \kappa/M_G^2$ is determined by instanton effects resulting from the new strong interaction with a suppression energy scale M_G. This is similar to the instanton effects found in Topcolor theories in Ref. [13]. In our calculations we will take κ to be a free parameter varying in the region $\kappa \simeq (0.1 - 1)$.

3 The Scalar Mass Spectrum

With the inclusion of the instanton-generated mixing term the most general potential that is generated by the fermion loops is given by

$$\begin{aligned}V(\Phi_U,\Phi_D) &= \mu_U^2|\Phi_U|^2 + \mu_D^2|\Phi_D|^2 + \mu_{UD}^2(\Phi_U^\dagger\Phi_D + {\rm h.c.}) \\ &\quad + \frac{\lambda_1}{2}|\Phi_U|^4 + \frac{\lambda_2}{2}|\Phi_D|^4 + \lambda_3|\Phi_U|^2|\Phi_D|^2 + \lambda_4|\Phi_U^\dagger\Phi_D|^2 \ ,\end{aligned} \tag{3.1}$$

where the scalar self-couplings λ_i, with $i = (1-4)$ are generated by fermion loops. The potential in (3.1) is not the most general 2HDM potential [14]. In general there would be other self-coupling terms allowed by the symmetries. However, these are not generated at one-loop level by the Yukawa interactions in (2.5).

The propagation of the scalars at low energies renormalizes the Yukawa couplings, leading to a renormalization group evolution given by

$$\frac{dY_U}{dt} = \frac{1}{16\pi^2}\left[Y_U^3 N_c + \frac{3}{2}Y_U^3 + \frac{1}{2}Y_D^2 Y_U - C_U(t)Y_U\right] \tag{3.2}$$

$$\frac{dY_D}{dt} = \frac{1}{16\pi^2}\left[Y_D^3 N_c + \frac{3}{2}Y_D^3 + \frac{1}{2}Y_U^2 Y_D - C_D(t)Y_D\right] \ , \tag{3.3}$$

where $t = \ln(\mu)$, and the functions $C_U(t)$and $C_D(t)$ give the contributions from the SM gauge sector. They are given by

$$\begin{aligned}C_U(t) &= 8g_s^2(t) + \frac{9}{4}g^2 + \frac{17}{12}g'^2 \\ C_D(t) &= 8g_s^2(t) + \frac{9}{4}g^2 + \frac{5}{12}g'^2 \ ,\end{aligned} \tag{3.4}$$

with g_s, g and g' the $SU(3)_c$, $SU(2)_L$ and $U(1)_Y$ gauge couplings respectively. We will assume that the new strong interactions in (2.2) do not break isospin symmetry. Thus, the only isospin breaking is introduced by the $U(1)_Y$ hypercharge interaction in (3.4). We can then safely approximate

$$Y_U \simeq Y_D \tag{3.5}$$

which results in a nearly degenerate fourth-generation quark doublet. With this approximation, we can easily solve the RGEs for the Yukawa couplings, resulting in

$$Y_U(\mu) = Y_D(\mu)\sqrt{\frac{C(\mu)}{5\left[1-\left(\frac{\mu}{\Lambda}\right)^{2C^2(\mu)/16\pi^2}\right]}}\ , \tag{3.6}$$

where we defined $C_U(\mu) \simeq C_D(\mu) = C(\mu)$ by neglecting the small hypercharge contributions. For instance, for $\mu = M_Z$, C_U and C_D differ by less than 1%.

Next, we consider the RGEs for the renormalized scalar self-couplings. These are given by

$$\begin{aligned}
\frac{d\lambda_1}{dt} &= \frac{1}{16\pi^2}\left[12\lambda_1^2 + 4\lambda_3^2 + 4\lambda_3\lambda_4 + 2\lambda_4^2 - 3\lambda_1(3g^2+g'^2) + \frac{3}{2}g^4 + \frac{3}{4}(g^2+g'^2)^2\right. \\
&\quad \left. +12\lambda_1 Y_U^2 - 12Y_U^4\right] && (3.7)\\
\frac{d\lambda_2}{dt} &= \frac{1}{16\pi^2}\left[12\lambda_2^2 + 4\lambda_3^2 + 4\lambda_3\lambda_4 + 2\lambda_4^2 - 3\lambda_2(3g^2+g'^2) + \frac{3}{2}g^4 + \frac{3}{4}(g^2+g'^2)^2\right. \\
&\quad \left. +12\lambda_2 Y_D^2 - 12Y_D^4\right] && (3.8)\\
\frac{d\lambda_3}{dt} &= \frac{1}{16\pi^2}\left[(\lambda_1+\lambda_2)(6\lambda_3+2\lambda 4) + 4\lambda_3^2 + 2\lambda_4^2 - 3\lambda_3(3g^2+g'^2) + \frac{9}{4}g^4 + \frac{3}{4}g'^4 - \frac{3}{2}g^2g'^2\right. \\
&\quad \left. +6\lambda_3(Y_U^2+Y_D^2) - 12Y_D^2Y_U^2\right] && (3.9)\\
\frac{d\lambda_4}{dt} &= \frac{1}{16\pi^2}\left[2(\lambda_1+\lambda_2)\lambda_4 + 4(2\lambda_3+\lambda_4)\lambda_4 - 3\lambda_4(3g^2+g'^2) + 3g^2g'^2\right. \\
&\quad \left. +6\lambda_4(Y_U^2+Y_D^2) + 12Y_D^2Y_U^2\right]\ . && (3.10)
\end{aligned}$$

The solution to the RGEs would respect the approximate relations

$$\lambda_1 \simeq \lambda_2 \simeq \lambda_3 \simeq -\lambda_4\ , \tag{3.11}$$

These would be obtained if one naively computes the λ_i's using only the fermion loops induced by the Yukawa interactions (2.5). Although the RGEs receive additional contributions to the evolution, such as fermion loops renormalizing the scalar wave-functions, these do not greatly modify these relations. Incidentally, the relations (3.11) among the scalar self-couplings satisfy the stability condition for the potential (3.1), which require that $\lambda_1, \lambda_2 > 0$, with $\sqrt{\lambda_1\lambda_2} > -\lambda_3 - \lambda_4$. The solution for the λ_1 is shown in Figure 3, for various values of the cutoff mass scale M_G.

With the above solutions for the Yukawa and scalar self-couplings, we can readily compute the scalar mass spectrum. Defining the ratio of VEVs as

$$\tan\beta \equiv \frac{v_D}{v_U}\ , \tag{3.12}$$

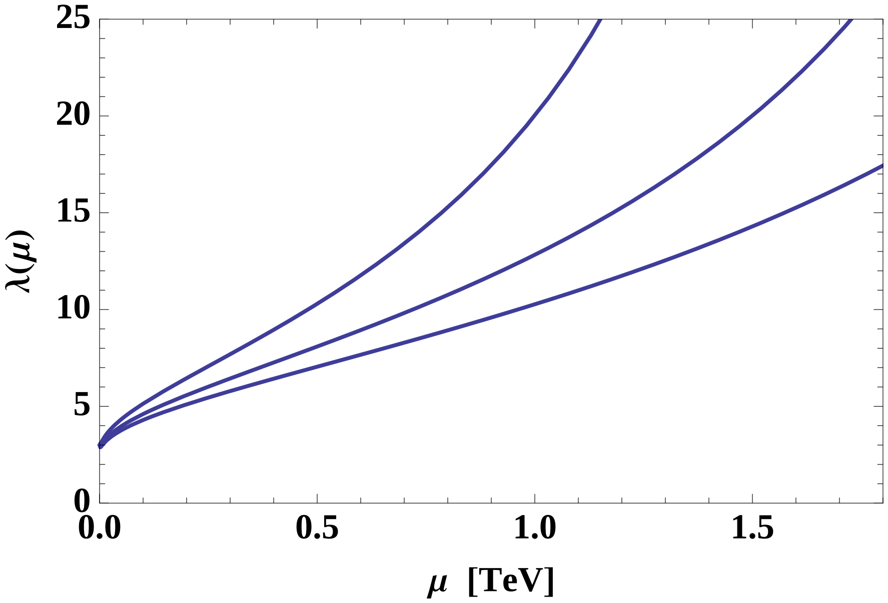

Figure 1. Scalar self-coupling λ_1 as a function of energy, for a various choices of the cutoff scale M_G. From the top $M_G = 2, 3, 4$ TeV.

we can express the scalar eigenvalues and eigenstates in terms of β and $v = \sqrt{v_U^2 + v_D^2}$.

In terms of the charged components of the Φ_U and Φ_D original doublets, the mass-eigenstate charged scalars are

$$H^{\pm} = \Phi_D^{\pm} \cos\beta - \Phi_U^{\pm} \sin\beta \ , \tag{3.13}$$

with masses given by

$$M_{H^{\pm}}^2 = -\frac{2\mu_{UD}^2}{\sin 2\beta} - \frac{\lambda_4}{2} v^2 \ . \tag{3.14}$$

The neutral states are

$$H = \sqrt{2}\left(Re[\Phi_U^0]\cos\gamma + Re[\Phi_D^0]\sin\gamma\right] \tag{3.15}$$
$$h = \sqrt{2}\left(-Re[\Phi_U^0]\sin\gamma + Re[\Phi_D^0]\cos\gamma\right] \ , \tag{3.16}$$

with masses

$$\begin{aligned} M_H^2 &= \left(\lambda_1 v^2 \cos^2\beta - \mu_{UD}^2 \tan\beta\right)\cos^2\gamma + \left(\lambda_2 v^2 \sin^2\beta - \mu_{UD}^2 \cot\beta\right)\sin^2\gamma \\ &\quad + \left[\mu_{UD}^2 + \frac{1}{2}(\lambda_3 + \lambda_4) v^2 \sin 2\beta\right] \sin 2\gamma \end{aligned} \tag{3.17}$$
$$\begin{aligned} M_h^2 &= \left(\lambda_1 v^2 \cos^2\beta - \mu_{UD}^2 \tan\beta\right)\sin^2\gamma + \left(\lambda_2 v^2 \sin^2\beta - \mu_{UD}^2 \cot\beta\right)\cos^2\gamma \\ &\quad - \left[\mu_{UD}^2 + \frac{1}{2}(\lambda_3 + \lambda_4) v^2 \sin 2\beta\right] \sin 2\gamma \ . \end{aligned} \tag{3.18}$$

The pseudo-scalar mass eigenstate defined by

$$A = \sqrt{2}\left(Im[\Phi_D^0]\cos\beta - Im[\Phi_U^0]\sin\beta\right) \ , \tag{3.19}$$

	$M_G = 2$ TeV	$M_G = 3$ TeV	$M_G = 4$ TeV
M_A	(26-118)	(15-59)	(10-39)
M_h	((548-580)	(459-467)	(422-425)
M_H	(651-732)	(530-537)	(482-585)
$M_{H^\pm}$	(603-719)	(495-512)	(453-459)

Table 1. The scalar spectrum for various values of the cutoff scale M_G. The ranges correspond to the instanton parameter range $k = (0.1 - 1)$ and are given in GeV.

has a mass is given by

$$M_A^2 = -2\frac{\mu_{UD}^2}{\sin 2\beta} \ . \tag{3.20}$$

The mixing angle γ is determined by the mass mixing induced by instantons, μ_{UD}, as well as from the terms in the potential involving λ_3 and λ_4 which result in $\Phi_U - \Phi_D$ mixing when two of the fields are replaced by their VEVs. Is given by

$$\tan 2\gamma = \frac{\mu_{UD}^2 + (\lambda_3 + \lambda_4)v^2 \sin 2\beta/2}{\mu_{UD}^2 + \lambda_4 v^2 \cos 2\beta/2} \ . \tag{3.21}$$

Finally, the renormalized up-down mixing mass parameter is given in terms of the instanton parameters as well as the renormalized self-couplings by

$$\mu_{UD}^2 = \frac{k\,v^2}{2M_G^2}\frac{\lambda_1\lambda_2\cos^2\beta\sin^2\beta}{\left[1 - kv^2(\lambda_1\cos^2\beta\cot\beta + \lambda_2\sin^2\beta\tan\beta)/(2M_G^2)\right]} \ . \tag{3.22}$$

We can now compute the scalar spectrum by defining a value for $\tan\beta$, the mass of the color octet M_G, and the order one parameter k from the instanton effects. In the context of the model presented here, it is clear that we will have

$$\tan\beta \simeq 1 \ , \tag{3.23}$$

given that we assume that the new strong interaction does not violate isospin. Furthermore, choosing M_G between $(2-4)$ TeV is reasonable to maintain naturalness in the dynamical mechanism to generate the electroweak scale. Finally, the instanton parameter is typically of order one: $k \simeq (0.1 - 1)$.

Within this range of parameters, the scalar spectrum varies little: typically we obtain scalar masses $M_{H_{1,2}} \simeq (500 - 600)$ GeV, and similarly for the charged scalar masses $M_{H^\pm}$. On the other hand, the pseudo-scalar mass M_A is much lighter in the range $M_A \simeq (20 - 200)$ GeV, and crucially depends on the instanton parameter k, as well as the suppression with the cutoff mass scale M_G. These results are summarized in Table 1. This is to be expected, since the scalar behaves as a pseudo-Nambu-Goldstone boson of the Peccei-Quinn symmetry, and its mass is only lifted by the explicit breaking due to instanton effects, which are dominated by the new interaction at the TeV scale. Thus, we see that the spectrum of the two Higgs doublet sector we obtain dynamically here is quite distinct and has all scalars in the 500 GeV range, with pseudo-scalar A being much lighter, perhaps as light as direct

bounds allow. This inverted spectrum has important consequences for the phenomenology of the scalar sector at colliders, as well as on the bounds, both direct and indirect from precision measurements. We comment on both of these in what follows.

4 Electroweak Precision Constraints

The presence of a new fermions results in additional contributions to electroweak precision observables. For illustrative purposes, we consider the case of a chiral fourth generation since it will involve also leptons, and it has a fixed number of degrees of freedom. The more general case is not necessarily as constraining on the parameter space. In addition to the fermionic contributions to precision observables, the formation of a composite two-Higgs doublet sector at low energies adds new contributions. The contributions to the S and T parameters from the two-Higgs doublet sector can be found, for instance, in Refs. [14] and [15]. The typical spectrum obtained by the double fourth-generation condensation, with a light pseudo-scalar and heavy and almost degenerate charged and CP-even scalars, results in a modest positive contribution to S. This is typically of the order of $S_{2H} \simeq 0.1$. On the other hand, the contributions of the scalar spectrum to T can vary more, approximately in the range $-0.2 < T_{2H} < 0.2$, for values of the cutoff in the range $\Lambda = (2-4)$ TeV , the instanton parameter taken in the interval $k = (0.1-1)$, and $\tan\beta = 0.9$.

These contributions must be added to the ones coming from the fourth-generation fermion loops [16, 17]. Since the composite scalar sector is made out of fourth-generation quarks, in principle these contributions to S and T are not completely independent. However, if we think in terms of a large-N_c expansion, the scalar contributions should be suppressed with respect to the single-quark loops. We will then add the two contributions coherently. The addition of these two contributions to S and T was done for much more general scalar and fermion spectra in Ref. [16]. In Figure 2 we plot the parameter space of the model considering the cutoff region $\Lambda = (2-4)$ TeV, the instanton parameter varying in the range $k = (0.1-1)$, and $\tan\beta = 0.9$. The RGE results from the previous section determine the overall scale of the fourth-generation quark spectrum. We take then the isospin splitting, which is mostly induced by electroweak corrections, to vary between $(0-100)$ GeV. In general, the overall mass scale of the lepton sector is not set by Λ, although it is related to it in some specific models [6]. To be general, we take neutrino masses to vary in the range $m_{\nu_4} = (100-500)\ GeV$, and again the isospin splitting with the charged lepton to be in the interval $(0-100)$GeV. The ellipses represent the 68% and 95% C.L. bounds from experiment as obtained in Ref. [18]. We can see that there is a region of parameter space of the model that is within the 95% C.L. interval. Thus, the scalar spectrum presented in the previous section, and the fourth-generation spectrum resulting from fourth-generation quark condensation are not excluded by electroweak data. The more general case, where the new fermions do not carry color, should be very similar and we take the case studied above as a concrete indication of the existence of a sizable paramater space.

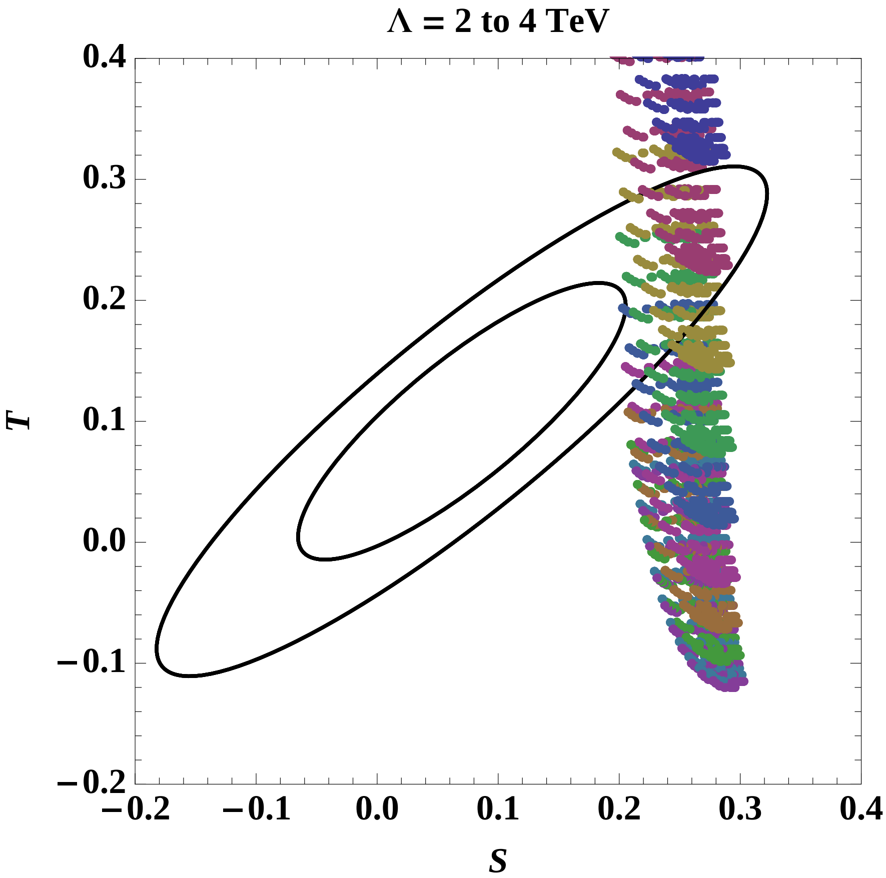

Figure 2. The contributions to the S and T parameters from a chiral fourth-generation and the composite two-Higgs sector resulting form quark condensation.

5 Phenomenology

In this section we consider some phenomenological aspects of these generic models. First, we address the issue of fermion masses and flavor violation. We then briefly discuss current bounds on and future searches for the scalar spectrum presented here.

5.1 Flavor Conservation and Constraints

Up to now, the model presented here breaks the electroweak symmetry and gives masses to the condensing fermions, e.g. the fourth-generation quarks. In order to give masses to all other fermions additional physics must be introduced. In general, the interaction giving rise to EWSB is not enough to generate the SM flavor structure. Higher dimensional operators such as

$$\bar{f}^i_L f^j_R \bar{f}^k_R f^\ell_L \;, \tag{5.1}$$

are needed. In (5.1), f^i denotes a fermion of generation i in a generic basis not necessarily aligned with the basis of the gauge interactions responsible for fourth-generation condensation. The size of the coefficients of these operators reflect a new dynamics, possibly at

a higher energy scale. For instance, in the spirit of the model of Ref. [6], there are four-quark operators generated by the exchange of the KK gluon. These become supercritically strong only for quark zero-modes highly localized in the IR. This is only the case for the fourth-generation. On the other hand, the flavor-violating operators in (5.1) are generated by non-renormalizable interactions in the 5D bulk, presumably resulting from a broken gauge flavor interaction at the UV scale. More generally, whatever the origin of the higher dimensional operators necessary to communicate EWSB to fermion masses, this can always be accommodated by having the doublet $\tilde{\Phi}_U$ couple only to up-type quarks, while having Φ_D couple only to down-type quarks and leptons. This is a two-Higgs model of type-II [14], and it does not result in tree-level neutral flavor violation.

In principle, we could worry about the fact that Φ_U and Φ_D are not mass eigen-states, since their mixing μ_{UD} is generated at low energies by instanton effects. However, as already pointed out in Ref. [11], the running of the dimensionless Yukawa couplings is unaffected by the mixing term, which only represents a soft breaking of the PQ symmetry. As a consequence, running to low energies preserves the fact that Φ_U and Φ_D only couple to up-type and down-type fermions respectively, ensuring that after mass diagonalization the neutral scalars have flavor-preserving neutral interactions.

Going beyond tree-level, the physical spectrum of the two-Higgs doublet model type II generates new contributions to various flavor observables [19]. The most stringent constraint comes from the contribution of the charged Higgs states to flavor-changing loop-induced decays, most notably $b \to s\gamma$. The data from B physics is consistent with the two-Higgs model spectrum as long as the 95% C.L. bound f $M_{H^\pm} > 316$ GeV is satisfied. In the present models, these states have masses well in excess of this bound, starting at about 500 GeV. On the other hand, the lighter pseudo-scalar state does not result in important contributions to loop-induced decays. We conclude that the two-Higgs doublet model spectrum favored by the condensation new chiral fermions described in Section 2 is not highly constrained by flavor data at the moment.

5.2 Direct Searches

The low energy bosonic spectrum of these models, with very heavy scalars $(h, H, H^{\pm})$ and a rather light pesudo-scalar A, presents a very distinct phenomenology. Although spectra of THDMs with A the lightest state have been studied in the context of supersymmetric models [20], as well as more generic THDM studies [21], the fact that the scalars are typically above 450 GeV presents important differences both for the existing bounds as well as the search strategies at the LHC. Although a detailed study of the phenomenology will be carried out elsewhere [22], there are some general remarks we can make here. First, existing bounds on M_A rely almost exclusively on its production in association with h at LEP. This bounds do not apply to our spectrum since $M_h > 400$ GeV means that LEP had no kinematic reach. In the case the new fermions are fourth-generation quarks, the bounds on the direct s-channel production of the scalars and pseudo-scalar are affected by the enhancement of the production cross sections due to the presence of the fourth-generation in the loop. This gives a cross section which is typically a factor of $(7-9)$ larger than that of a standard THDM type II with three generations. This is not the case for the colorless

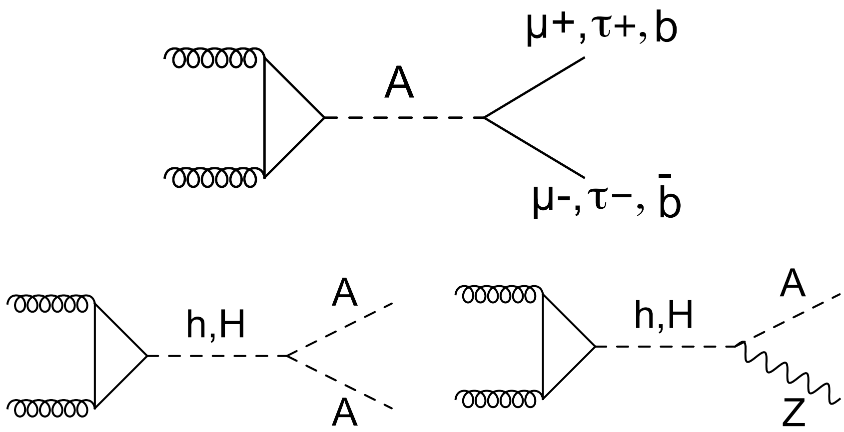

Figure 3. Diagrams dominating the production of the pseudo-scalar A at hadron colliders.

fermion models, where only the $\phi\gamma\gamma$ couplings are affected significantly. Considering the light end of the pseudo-scalar state, the bounds from B factories are $M_A > 9.5$ GeV at 90% C.L. Tevatron searches are not sensitive to this flipped spectrum. The LHC searches are beginning to be constraining in the presence of a fourth generation. However, and although the SM Higgs is excluded up to masses of $\simeq 600$ GeV in this case, this is yet not applicable to the two-Higgs doublet spectrum presented here.

As we can see from Table 3, M_A can be anywhere between 10 and 120 GeV, and possibly somewhat beyond both ends of this range. The s-channel A production gives the largest cross section, but it represents a challenge, especially at the larger masses where the branching ratio into b quarks dominates. Although the modes mediated by the heavy CP-even neutral scalars is suppressed due to their large masses, the enhancement of the cross section due to the additional fermions in the loop makes this modes of great interest [22], especially for models with additional colored doublets. The CP-even scalars (h, H), would be produced with a larger cross section too, but the addition of new important decay channels such as AA and AZ change the direct limits obtained for the SM Higgs decaying into W^+W^-, ZZ by ATLAS and CMS [24]. A full phenomenological study of the branching ratios and preferred decays for a given spectrum for both the colored and colorless fermion cases is left for a separate publication [22].

6 Conclusions and Outlook

We have shown that electroweak symmetry breaking by fermion condensation typically leads to a two-Higgs doublet model at low energies. We modeled the scalar spectrum of these scenarios by using the NJL approach. We obtained an inverted scalar spectrum, with the neutral CP-even and charged scalars heavy $(h, H, H^{\pm})$, typically at $(500 - 600)$ GeV, whereas the pseudo-scalar A is much lighter due to a Peccei-Quinn symmetry in the original action, only broken by the instantons of the new strong interactions.

The low energy scalar spectrum corresponds to a type-II two-Higgs doublet model. Thus tree-level flavor conservation is built in. For the masses obtained, loop-induced flavor-changing processes are not binding. mostly since $m_{h^{\pm}} > 450$ GeV avoids $b \to s\gamma$ bounds. The study of electroweak precision observables shows that per se the scalar spectrum does not violate S and T bounds significantly. However, when a fourth generation is included important regions of parameter space are excluded. This exclusion is independent of whether the new fermion sector is a fourth generation or a more exotic matter content, since the contributions to S and T depend only on the number of fermions, and their mass splittings. However, we conclude that an important region of the parameter space is still allowed by electroweak precision bounds, as can be seen from Figure 2.

This inverted two-Higgs doublet model spectrum leads to a distinct phenomenology at the LHC, with the typical decay channel for the neutral CP even states, $(h, H) \to WW, ZZ$ now being in competition with $(h, H) \to AA, AZ$. These new channels, involving either 4 b's, τ's or 2 b's and 2 charged leptons, should be looked for at the LHC, for the appropriate regions of the model's parameter space. We expect that the LHC with $\sqrt{s} = 7$ TeV will be able to fully test these regions. A detailed study of this phenomenology will be undertaken in Ref. [22].

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